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Constraints on left-right symmetric models from the process $b \rightarrow s\gamma$

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ABSTRACT

In left-right symmetric models, large contributions to the decay amplitude $b \rightarrow s\gamma$ can arise from the mixing of the W_L and W_R gauge bosons as well as from the charged Higgs boson. These amplitudes are enhanced by the factor m_t/m_b compared to the contributions in the standard model. We use the recent CLEO results on the radiative B decay to place constraints on the $W_L - W_R$ mixing angle ζ and the mass of the charged Higgs boson m_{H^\pm} . Significant departures from the standard model predictions occur when $|\zeta| \gtrsim 0.003$ and/or when $m_{H^\pm} \lesssim$ a few TeV.

Introduction:

Left-right symmetric theories of the weak interactions based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_Y$ are attractive extensions of the standard model possessing manifest parity invariance [1]. These theories also have greater quark-lepton symmetry than the standard model since they require the existence of the right-handed partner of the neutrino ν_R , leading naturally to non-zero neutrino masses. The observed (V-A) nature of the weak interactions is explained by the spontaneous breaking of parity along with the breaking of $SU(2)_R$ gauge symmetry at a scale $v_R \gg m_W$. If the scale v_R of $SU(2)_R$ breaking is not much above the weak scale, observable deviations from the predictions of the standard model are possible. Flavor changing neutral current processes have proven in the past to be powerful probes of physics beyond the standard model. For example, in the context of the left-right symmetric models, the mass of the charged W_R gauge boson should exceed about 1.6 TeV, or else it would contribute to the $K^0 - \bar{K}^0$ mass difference at an unacceptable level [2].

In this paper we study the constraints on the parameters of the left-right symmetric model arising from the process $b \rightarrow s\gamma$. Recently the CLEO collaboration has reported the first observation of the exclusive decay $B \rightarrow K^*\gamma$ with a branching ratio [3]

$$Br(B \rightarrow K^*\gamma) = (4.5 \pm 1.5 \pm 0.9) \times 10^{-5}. \quad (1)$$

Eq. (1) implies both lower and an upper limits on the inclusive decay $B \rightarrow X_s\gamma$ [e.g. $Br(B \rightarrow X_s\gamma) < 5.4 \times 10^{-5}$ at 95% C.L.]. These numbers are in good agreement with the standard model predictions and as such, are sensitive to new physics.

In renormalizable gauge theories, the radiative decay $b \rightarrow s\gamma$ proceeds through the magnetic moment operators $\bar{b}_R \sigma_{\mu\nu} s_L F^{\mu\nu}$ and $\bar{b}_L \sigma_{\mu\nu} s_R F^{\mu\nu}$, where $F^{\mu\nu}$ is the electromagnetic field strength tensor. In the standard model, the $b \rightarrow s\gamma$ amplitude is proportional to m_b or m_s , the mass of the bottom quark or the strange quark, because the pure (V-A) structure of the charged currents requires the chirality-flip to proceed only through the mass of the initial or the final state quark. In contrast, in left-right symmetric models, the mixing of the W_L and W_R gauge bosons leads also to (V+A) interactions between the W_1 boson and the quarks, where W_1 is the lighter mass eigenstate formed by W_L

and W_R . In this case, the $b \rightarrow s\gamma$ amplitude can be proportional to the top quark mass m_t rather than m_b or m_s since chirality flip can now occur with the top quark mass in the intermediate state. This enhancement of the amplitude gives rise to significant departure of the decay rate $Br(b \rightarrow s\gamma)$ from the prediction in the standard model, if the $W_L - W_R$ mixing angle ζ exceeds about 10^{-3} .

Left-right symmetric models also predict the existence of a charged Higgs boson that couples to the quarks. Its contributions to the $b \rightarrow s\gamma$ amplitude are also proportional to the top quark mass, and the experimental result (1) already probes the charged Higgs mass of a few TeV. This feature should be compared to the charged Higgs contributions to the $b \rightarrow s\gamma$ amplitude in the minimal supersymmetric standard model (MSSM) which are proportional to m_b or m_s . The experimental result (1) excludes a charged Higgs boson lighter than a few hundred GeV in this case [4]. The enhancement of the charged Higgs boson contributions in left-right symmetric models stems from the absence of natural flavor conservation in the Higgs sector of the model. In spite of the absence of flavor conservation in left-right symmetric models, the interactions of the charged Higgs boson to the quarks are determined in terms of the quark masses, Cabibbo-Kobayashi-Maskawa (CKM) mixing angles and the ratio of vacuum expectation values, just as in the MSSM.

Radiative b -decays have been studied in the context of left-right symmetric models in the past in Refs. [5, 6]. The effects of the $W_L - W_R$ mixing on the $b \rightarrow s\gamma$ amplitude were studied in Ref. [5], but the contributions from the charged Higgs boson were not examined there. Moreover, our result on the contributions of the $W_L - W_R$ mixing disagrees with that in Ref. [5]. The charged Higgs contributions were analyzed in Ref. [6], but the effects of the $W_L - W_R$ mixing and the leading QCD corrections were not included. In realistic left-right models, large contributions to $b \rightarrow s\gamma$ amplitude arising from the $W_L - W_R$ mixing and those from the charged Higgs boson are closely related to each other and they can be simultaneously sizable. Here we present a comprehensive analysis of both of these contributions taking into account their correlations and clarify the implications of the recent $b \rightarrow s\gamma$ experiment on the parameters in the left-right symmetric models.

Left-right symmetric models:

Left-right symmetric models of weak interactions are based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The quarks (q) and leptons (l) transform under the gauge group as

$$\begin{aligned} q_L(2, 1, \frac{1}{3}) &= \begin{pmatrix} u \\ d \end{pmatrix}_L ; \quad q_R(1, 2, \frac{1}{3}) = \begin{pmatrix} u \\ d \end{pmatrix}_R \\ l_L(2, 1, -1) &= \begin{pmatrix} \nu \\ e \end{pmatrix}_L ; \quad l_R(1, 2, -1) = \begin{pmatrix} \nu \\ e \end{pmatrix}_R \end{aligned} \quad (2)$$

where generation indices have been suppressed. The minimal Higgs sector compatible with the see-saw mechanism for small neutrino masses consists of the multiplets [7] $\Delta_L(3, 1, 2)$, $\Delta_R(1, 3, 2)$ and $\Phi(2, 2, 0)$ which in component form read as

$$\Delta_{L,R} = \begin{pmatrix} \delta^+/\sqrt{2} & \delta^{++} \\ \delta^0 & -\delta^+/\sqrt{2} \end{pmatrix}_{L,R}, \quad \Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix}. \quad (3)$$

The field Δ_R is needed for breaking the gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ down to the gauge symmetry in the standard model and to give Majorana masses to the right-handed neutrinos. The field Φ is required for generating the quark and lepton masses. The field Δ_L is present in the theory to maintain the discrete parity invariance. Under parity transformation, $q_L \rightarrow q_R$, $l_L \rightarrow l_R$, $\Delta_L \rightarrow \Delta_R$, $\Phi \rightarrow \Phi^\dagger$ and $W_L \rightarrow W_R$.

Spontaneous breaking of the gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ down to $U(1)_{EM}$ is achieved by the vacuum expectation values (VEV) of the neutral Higgs fields denoted by

$$\langle \Delta_{L,R} \rangle = \begin{pmatrix} 0 & 0 \\ v_{L,R} & 0 \end{pmatrix}, \quad \langle \Phi \rangle = \begin{pmatrix} k & 0 \\ 0 & k' \end{pmatrix}. \quad (4)$$

Among the vacuum expectation values k , k' and $v_{L,R}$, the hierarchy $k, k' \ll v_R$ is needed to preserve the success of the standard (V-A) theory. In this case, another hierarchy $v_L \ll k, k'$ follows from a detailed analysis of the Higgs potential [7] which yields the relation $v_L \sim \gamma k^2/v_R$, where γ is some combination of the Higgs quartic coupling constants. This is a welcome result since the analysis of the electroweak ρ -parameter leads to the constraint $v_L \lesssim 10$ GeV [8] and a natural realization of the see-saw mechanism for small neutrino masses requires $v_L \lesssim$ a few MeV. In what follows, we shall work in the limit $v_L \rightarrow 0$, which is justified for the above reasons. The VEVs k or k' can in general have a

phase, but we shall assume this phase to be small. This is also justified from the detailed analysis of the Higgs potential [9]. Small non-zero values of v_L or the relative phase will not alter our conclusions.

The Yukawa Lagrangian involving the quark fields is given by

$$\mathcal{L}_Y = \bar{q}_L h \Phi q_R + \bar{q}_L \tilde{h} \tilde{\Phi} q_R + h.c., \quad (5)$$

where $\tilde{\Phi} \equiv \tau_2 \Phi^* \tau_2$, h and \tilde{h} are 3×3 hermitian matrices in generation space. Eq. (5) leads to the following mass matrices for the up-type and down-type quarks:

$$M_u = hk + \tilde{h}k', \quad M_d = hk' + \tilde{h}k. \quad (6)$$

In the charged gauge boson sector, the W_L^\pm and W_R^\pm mix with their mass-squared matrix given by

$$\mathcal{M}^2 = \frac{g^2}{2} \begin{pmatrix} k^2 + k'^2 & -2kk' \\ -2kk' & 2v_R^2 + k^2 + k'^2 \end{pmatrix}. \quad (7)$$

The two mass eigenstates are

$$\begin{aligned} W_1^\pm &= c_\zeta W_L^\pm + s_\zeta W_R^\pm, \\ W_2^\pm &= -s_\zeta W_L^\pm + c_\zeta W_R^\pm, \end{aligned} \quad (8)$$

where $s_\zeta = \sin\zeta$, $c_\zeta = \cos\zeta$, respectively, and

$$\tan 2\zeta = \frac{2kk'}{v_R^2}. \quad (9)$$

We have defined $m_{W_1} \leq m_{W_2}$ with $m_{W_1} \simeq 80$ GeV. The mass eigenvalues of W_1^\pm and W_2^\pm are given by

$$m_{W_1}^2 = \frac{g^2}{2}(k^2 + k'^2 - 4kk's_\zeta c_\zeta + 2v_R^2 s_\zeta^2), \quad m_{W_2}^2 = \frac{g^2}{2}(k^2 + k'^2 + 4kk's_\zeta c_\zeta + 2v_R^2 c_\zeta^2). \quad (10)$$

The coupling of the lighter charged W_1 -boson to the quarks is given by

$$\mathcal{L}_{W_1} = \frac{g}{2\sqrt{2}} \left(\bar{u}, \bar{c}, \bar{t} \right) \{c_\zeta \gamma^\mu (1 - \gamma_5) + s_\zeta \gamma^\mu (1 + \gamma_5)\} W_{1\mu}^+ V \begin{pmatrix} d \\ s \\ b \end{pmatrix} + h.c., \quad (11)$$

where V is the Cabibbo-Kobayashi-Maskawa (CKM) matrix. Note that since the up and down mass matrices are hermitian (as the VEVs k, k' are taken to be real), the right-handed CKM matrix and the left-handed CKM matrix are equal ($V_L = V_R = V$), which is reflected in Eq. (11). The interaction (5) leads to the following coupling of the corresponding (unphysical) Nambu-Goldstone boson G_1 to the quarks,

$$\begin{aligned} \mathcal{L}_{G_1} = & \frac{g}{2\sqrt{2}m_{W_1}} \left(\bar{u}, \bar{c}, \bar{t} \right) c_\zeta \{(1 - \gamma_5)D_u V - (1 + \gamma_5)V D_d\} G_1^+ \begin{pmatrix} d \\ s \\ b \end{pmatrix} \\ & + \frac{g}{2\sqrt{2}m_{W_1}} \left(\bar{u}, \bar{c}, \bar{t} \right) s_\zeta \{(1 + \gamma_5)D_u V - (1 - \gamma_5)V D_d\} G_1^+ \begin{pmatrix} d \\ s \\ b \end{pmatrix} + h.c., \quad (12) \end{aligned}$$

where we define the diagonal mass matrices $D_u = \text{diag}(m_u, m_c, m_t)$ and $D_d = \text{diag}(m_d, m_s, m_b)$. The couplings of the heavier W_2^\pm gauge boson and the analogous Nambu-Goldstone boson G_2 can be obtained from Eqs. (11)-(12) by the replacement $s_\zeta \rightarrow c_\zeta, c_\zeta \rightarrow -s_\zeta$ and $\gamma_5 \rightarrow -\gamma_5$, but are not necessary for the present analysis.

The model has two physical charged Higgs bosons. In the limit of $v_L \rightarrow 0$, one of them, δ_L^\pm , becomes mass eigenstate by itself. The Higgs boson δ_L^\pm has couplings only to the leptons, and does not enter into the discussion of the $b \rightarrow s\gamma$ amplitude. The second physical Higgs boson, H^\pm , which is the linear combination orthogonal to $G_{1,2}^\pm$,

$$H^\pm = N_{H^\pm} \left[k' \phi_1^\pm + k \phi_2^\pm + \frac{(k^2 - k'^2)}{\sqrt{2}v_R} \delta_R^\pm \right], \quad (13)$$

has the following Yukawa coupling to the quarks:

$$\begin{aligned} \mathcal{L}_{H^\pm} = & -\frac{\sin(2\beta)N_{H^\pm}}{2\cos(2\beta)} \left(\bar{u}, \bar{c}, \bar{t} \right) \{(1 - \gamma_5)D_u V - (1 + \gamma_5)V D_d\} H^\pm \begin{pmatrix} d \\ s \\ b \end{pmatrix} \\ & - \frac{N_{H^\pm}}{2\cos(2\beta)} \left(\bar{u}, \bar{c}, \bar{t} \right) \{(1 + \gamma_5)D_u V - (1 - \gamma_5)V D_d\} H^\pm \begin{pmatrix} d \\ s \\ b \end{pmatrix} + h.c., \quad (14) \end{aligned}$$

where

$$N_{H^\pm} = 1/\sqrt{k^2 + k'^2 + \frac{(k^2 - k'^2)^2}{2v_R^2}}, \quad \tan\beta = k/k'. \quad (15)$$

The mass of this charged Higgs boson is dependent on the detailed structure of the Higgs potential, and we leave it as a free parameter in our analysis.

The process $b \rightarrow s\gamma$:

We now investigate the effective Hamiltonian describing $b \rightarrow s\gamma$ in the left-right symmetric model. Using the Lagrangians (11), (12) and (14), the effective Hamiltonian for $b \rightarrow s\gamma$ decay can be written as

$$H_{eff} = \frac{e}{16\pi^2} \frac{2G_F}{\sqrt{2}} V_{tb} V_{ts}^* m_b (A_L \bar{s}_L \sigma^{\mu\nu} b_R + A_R \bar{s}_R \sigma^{\mu\nu} b_L) F_{\mu\nu}, \quad (16)$$

with

$$A_L = A_{SM}(x) + \zeta \frac{m_t}{m_b} A_{RH}(x) + \frac{m_t s_{2\beta}}{m_b c_{2\beta}^2} A_{H^+}^1(y) + \tan^2(2\beta) A_{H^+}^2(y), \quad (17)$$

$$A_R = \zeta \frac{m_t}{m_b} A_{RH}(x) + \frac{m_t s_{2\beta}}{m_b c_{2\beta}^2} A_{H^+}^1(y) + \frac{1}{c_{2\beta}^2} A_{H^+}^2(y). \quad (18)$$

Here the masses of the light quarks u, d, s and c have been neglected and the approximations $c_\zeta \simeq 1$, $s_\zeta \simeq \zeta$ and $N_{H^+} \simeq g/(\sqrt{2}m_W)$ have been used. In eq. (16), $F_{\mu\nu}$ is the electromagnetic field strength tensor, $x = m_t^2/m_W^2$, $y = m_t^2/m_{H^+}^2$, $s_{2\beta} = \sin 2\beta$ and $c_{2\beta} = \cos 2\beta$. The functions A_{SM} , A_{RH} , $A_{H^+}^1(y)$ and $A_{H^+}^2(y)$ are found to be

$$\begin{aligned} A_{SM}(x) &= \frac{1}{(1-x)^4} Q_t \left\{ \frac{x^4}{4} - \frac{3}{2}x^3 + \frac{3}{4}x^2 + \frac{x}{2} + \frac{3}{2}x^2 \log(x) \right\} \\ &\quad + \frac{1}{(1-x)^4} \left\{ \frac{x^4}{2} + \frac{3}{4}x^3 - \frac{3}{2}x^2 + \frac{x}{4} - \frac{3}{2}x^3 \log(x) \right\}, \\ A_{RH}(x) &= \frac{1}{(1-x)^3} Q_t \left\{ -\frac{x^3}{2} - \frac{3}{2}x + 2 + 3x \log(x) \right\} \\ &\quad + \frac{1}{(1-x)^3} \left\{ -\frac{x^3}{2} + 6x^2 - \frac{15}{2}x + 2 - 3x^2 \log(x) \right\}, \\ A_{H^+}^1(y) &= \frac{1}{(1-y)^3} Q_t \left\{ -\frac{y^3}{2} + 2y^2 - \frac{3}{2}y - y \log(y) \right\} + \frac{1}{(1-y)^3} \left\{ -\frac{y^3}{2} + \frac{y}{2} + y^2 \log(y) \right\}, \\ A_{H^+}^2(y) &= \frac{1}{3} A_{SM}(y) - A_{H^+}^1(y), \end{aligned} \quad (19)$$

where $Q_t = 2/3$ is the electric charge of the top-quark.

A_{SM} in eq. (17) is the contribution given by the standard model [10]. The right-handed coupling in eq. (11) leads to the contributions $(m_t/m_b)\zeta A_{RH}$ [11] in eqs. (17) and (18). These contributions arise from the chirality flip induced by the top quark mass in the intermediate state, and they are enhanced by the factor m_t/m_b compared to the

standard model. Chirality-flip inside the loop is forbidden in the standard model because of the purely left-handed nature of the W-boson coupling to the quarks. The last two terms in eqs. (17) and (18) are the contributions from the charged Higgs boson H^+ . Note that these are also enhanced by the factor m_t/m_b compared to the contributions in the standard model. This is in contrast with the charged Higgs contributions in the minimal supersymmetric standard model (MSSM), which are proportional to m_b or m_s , without any large enhancement factor [4]. It is worthwhile to emphasize the correlation between the $W_L - W_R$ mixing contributions and the charged Higgs contributions in eqs. (17) and (18). From the expressions (9) and (10), the mixing angle ζ can be written as $\zeta \simeq \sin(2\beta) \cdot (m_{W_1}^2/m_{W_2}^2)$. Therefore the large contributions coming from the $W_L - W_R$ mixing and the charged Higgs boson are proportional to each other and they can be sizable if k and k' are of the same order.

The effective Hamiltonian H_{eff} given in eq. (16) has been evaluated at the electroweak scale ($\mu \sim m_W$). To make contact with the $b \rightarrow s\gamma$ decay, H_{eff} should be evolved down to lower momentum scale ($\mu \sim m_b$) by the renormalization group analysis. The leading QCD corrections to the Hamiltonian (16) during its evolution turn out to be significant [12, 13]. These QCD corrections have been computed in the standard model in Refs. [14, 15] by analyzing the operator mixing between the magnetic moment operator $\bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu}$ in eq. (16) and the four Fermi operators involving the quarks lighter than the W bosons. In left-right models, because of the $W_L - W_R$ mixing, there exists some new four-Fermi operators which mix with the magnetic moment operator $\bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu}$. However, the effects of these new operators are simply order ζ *without the enhancement factor of m_t/m_b* , and are negligible in our analysis. The running of the strong coupling constants and the effects of the operator mixing are also negligible in the momentum region above the W_1 boson mass, because of the asymptotic freedom of the strong interactions, and consequently the effects of new scales characterized by m_t and m_{H^+} are ignored. The QCD corrections to the other magnetic moment operator $\bar{s}_R \sigma^{\mu\nu} b_L F_{\mu\nu}$ can be computed in analogy to the case of the operator $\bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu}$ because the strong interactions respect parity. Therefore, we compute the QCD corrections to the Hamiltonian (16) following the procedure of Ref. [14, 15] established in the standard model. Here we use the simplified analytical results which are exact to within a few

percent [14]. Including the QCD corrections, A_L and A_R in the effective Hamiltonian (16) are renormalized as

$$\begin{aligned} A_L^{eff} &= \eta^{-32/23} \{A_L + \frac{3}{10} X(\eta^{10/23} - 1) + \frac{3}{28} X(\eta^{28/23} - 1)\}, \\ A_R^{eff} &= \eta^{-32/23} A_R. \end{aligned} \quad (20)$$

with $X = 208/81$ [15] and $\eta = \alpha_s(m_b^2)/\alpha_s(m_W^2) \simeq 1.8$. The last two terms in A_L^{eff} come from the operator mixing in the standard model. Analogous contributions to A_R^{eff} in the standard model are proportional to the mass of the strange quarks and are neglected in our analysis.

The branching fraction $Br(b \rightarrow s\gamma)$ is computed following the procedure of Ref. [12] by normalizing the decay width $\Gamma(b \rightarrow s\gamma)$ to the semileptonic decay width $\Gamma(b \rightarrow ce\bar{\nu})$,

$$Br(b \rightarrow s\gamma) = \frac{\Gamma(b \rightarrow s\gamma)}{\Gamma(b \rightarrow ce\bar{\nu})} Br(b \rightarrow ce\bar{\nu}), \quad (21)$$

and using $Br(b \rightarrow ce\bar{\nu}) \simeq 11\%$ [16]. Using the Hamiltonian (16) with renormalized quantities $A_{L,R}^{eff}$, the rate for $\Gamma(b \rightarrow s\gamma)$ normalized to the semileptonic rate is given by,

$$\frac{\Gamma(b \rightarrow s\gamma)}{\Gamma(b \rightarrow ce\bar{\nu})} = \frac{3\alpha}{2\pi\rho(m_c/m_b)(1 - \delta_{QCD})} (|A_L^{eff}|^2 + |A_R^{eff}|^2) \quad (22)$$

In eq. (22), $\rho(m_c/m_b)$ and δ_{QCD} are the phase space suppression factor and the QCD corrections to the semileptonic decay, respectively. These factors are evaluated as $\rho(m_c/m_b) = 0.447$ and $\delta_{QCD} = (2\alpha_s(m_b^2)/3\pi)f(m_c/m_b)$, where $f(m_c/m_b) = 2.41$ [17]. (We take $\alpha_s(m_b^2) = 0.23$.)

We investigate implications of eqs. (21) and (22) numerically. In Fig. 1, we plot the branching ratio as a function of the left-right mixing angle ζ for three different values of the top-quark mass, $m_t = 110$ GeV, 140 GeV and 170 GeV. The value of m_b appearing in the enhancement factor m_t/m_b in eqs. (17) and (18) is the one evaluated at $\mu \sim m_t$. We choose it to be 3 GeV corresponding to a pole mass of 4.8 GeV. Here we have kept the charged Higgs boson mass m_{H^\pm} rather high ($m_{H^\pm} = 20$ TeV). In this case the contributions of the charged Higgs boson are negligible. The CLEO limit on the branching ratio, $Br(b \rightarrow s\gamma) \leq 5.4 \times 10^{-5}$ implies a limit $-0.015 \leq \zeta \leq 0.003$, with the region around $\zeta = -0.005$ disfavored. This limit should be compared to the

existing bounds on ζ in left-right models. A limit $|\zeta| \leq 0.035$ has been inferred from the measurement of the ξ -parameter in polarized μ decay [16], but it assumes the neutrino to be a Dirac particle, which is not the case in the popular see-saw mechanism. A bound $|\zeta| \leq 0.004$ has been derived from non-leptonic K decays using the MIT bag model and assuming current algebra and PCAC [18], but this bound is clouded by traditional strong interaction uncertainties. The bound derived here from the $b \rightarrow s\gamma$ process holds regardless of the nature of the neutrino and has considerably less strong interaction uncertainty.

In Fig. 2, we have plotted the branching ratio as a function of $\sin(2\beta)$ for various Higgs masses ($m_{H^\pm} = 0.5, 1, 3, 5$ and 10 TeV) with $m_t = 140$ GeV kept fixed. Here we have set $\zeta = 0$, so that the entire non-standard contribution arises from the charged Higgs sector. When $\sin(2\beta)$ approaches ± 1 (keeping the CKM matrix elements fixed), the coupling of the charged Higgs boson to the quarks diverges. By requiring that these couplings should be perturbatively controllable, i.e., less than 4π , we find that $|s_{2\beta}| \lesssim 0.98$. Charged Higgs boson with the mass below a few TeV will contradict the CLEO results on $b \rightarrow s\gamma$ if k and k' are of the same order.

In Fig. 3, we plot contours of branching ratios for various Higgs boson masses ($m_{H^\pm} = 0.5, 1, 3, 5, 10$ TeV) for non-zero values of ζ and $m_t = 140$ GeV. We have fixed $m_{W_2} = 1.6$ TeV for this graph (which satisfies the indirect bound from $K^0 - \bar{K}^0$ mass difference) and used the relation $\zeta \simeq \sin(2\beta) \cdot m_{W_1}^2 / m_{W_2}^2$ to determine ζ for a given $\sin(2\beta)$. The contributions from the mixing angle ζ and the charged Higgs boson *act additively* to the effective Hamiltonian (16). Consequently, the charged Higgs boson lighter than a few TeV are excluded for a wide range of parameter space even after the inclusion of the effects of ζ . In particular, the charged Higgs boson with a mass $m_{H^+} \lesssim 1$ TeV is allowed only if k and k' differ by an order of magnitude or so. In Fig. 4 we plot analogous contours, but corresponding to $m_{W_2} = 800$ GeV. A wider range of parameters are excluded by $b \rightarrow s\gamma$ in this case.

Conclusions:

We have examined the decay $b \rightarrow s\gamma$ in left-right symmetric models. Large contribution to $b \rightarrow s\gamma$ amplitude arises from $W_L - W_R$ mixing. The physical charged Higgs boson present in the minimal Higgs sector of the model also yields significant contribution to the decay amplitude. Both these contributions are enhanced by the factor m_t/m_b compared to the standard model or the minimal supersymmetric standard model. This enhancement stems from the chirality-flip induced by the right-handed coupling of the W_1 -boson to the quarks in the case of the the $W_L - W_R$ mixing. In the case of the charged Higgs contribution, the enhancement is closely related to the absence of flavor conservation in the Higgs sector. Because of these enhanced contributions, the decay $b \rightarrow s\gamma$ can serve as a sensitive probe to possible signals of left-right symmetric models. The recent CLEO results on the radiative B decay lead to the most stringent and essentially model-independent bound on the $W_L - W_R$ mixing angle ζ in a general class of left-right models: $-0.015 \leq \zeta \leq 0.003$. The mass of the charged Higgs boson, which has not been probed so far by other experiments, is also stringently constrained by the $b \rightarrow s\gamma$ experiments. In particular, the charged Higgs boson mass lighter than about a few TeV is excluded for a wide range of parameter space.

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Note added

While preparing this manuscript, we received a preprint by P. Cho and M. Misiak (CALT-68-1893, hep-ph 9310332). They examine the effect of the $W_L - W_R$ mixing on the $b \rightarrow s\gamma$ amplitude, however, the charged Higgs contributions are not considered there. Our result on the contributions of the $W_L - W_R$ mixing are in agreement with theirs. We

also found that in the calculation of the QCD corrections to the effective Hamiltonian (16), the effects of the new operators ($O_{9,10}$ in their paper) on the analysis of the operator mixing are small due to the reasons we explained in the text. We would like to thank M. Misiak for a clarifying discussion on the QCD corrections.

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Figure Captions

Fig. 1 The branching fraction $Br(b \rightarrow s\gamma)$ in the left-right symmetric model as a function of the $W_L^+ - W_R^+$ mixing angle ζ . The three curves correspond to $m_t=110$ GeV, 140 GeV, and 170 GeV. The charged Higgs boson mass is set equal to 20 TeV so that its effects are negligible for these curves.

Fig. 2 The branching ratio $Br(b \rightarrow s\gamma)$ as a function of the mixing angle $\sin(2\beta)$ for various values of the charged Higgs boson mass. $m_{H^+} = 0.5$ TeV (inner solid), 1 TeV (inner dotdash), 3 TeV (dash), 5 TeV (outer dotdash) and 10 TeV (outer solid). The top mass is fixed as $m_t=140$ GeV. This graph corresponds to $\zeta = 0$.

Fig. 3 $Br(b \rightarrow s\gamma)$ versus $\sin(2\beta)$ for $m_t=140$ GeV and $m_{W_2}=1.6$ TeV for the same set of m_{H^+} values as in Fig. 2.

Fig. 4 Same as in Fig. 3, but for $m_{W_2} = 800$ GeV.

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